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Fermion Masses without Higgs: A Supersymmetric Technicolor Model

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Abstract

We propose a supersymmetric technicolor model in which the electroweak symmetry breaking is communicated to the quarks and leptons by technicolored $SU(2)_W$-singlet scalars. When the technifermions condense, the quarks and leptons of the third generation acquire mass. The fermions of the other generations do not couple to the technicolored scalars but they receive masses from radiative corrections involving superpartners. As a result, the mass hierarchy between the fermion generations arises naturally. The model predicts the CP asymmetries in B meson decays and in $\Delta S = 1$ transitions to be smaller by two orders of magnitude than the ones predicted in the Standard Model.

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1 Introduction

The Higgs doublet has a double role in the Standard Model (SM): to break the electroweak symmetry spontaneously and to give mass to the fermions. The latter offers no explanation for the pattern of masses of the quarks and leptons. The mixing angles of the quarks and the CP violation phase are also free parameters in the SM. Moreover, the arbitrariness in the phase of the quark mass determinant is one of the sources of the strong CP problem [1].

The existence of the light Higgs boson in the SM is unnatural [2] because of the quadratic divergences in the scalar self-energy. This naturalness problem can be solved while maintaining fundamental scalars in a theory with supersymmetry (SUSY) broken softly [3]. However, the simple structure of the Higgs sector of the SM is lost in the Supersymmetric Standard Model (SSM) where there is need for two Higgs doublets to provide mass for both up-type and down-type quarks. Furthermore, in the SSM the constraints from flavor-changing neutral currents (FCNC) require a high degeneracy between the squarks with the same charge [4], which could occur only if strong assumptions are imposed [5, 6, 7].

In technicolor models [8] the electroweak symmetry is broken dynamically and there is no need for a Higgs doublet provided a mechanism for fermion mass generation is found. In extended technicolor (ETC) [9] the $SU(2)_W \times U(1)_Y$ symmetry breaking is transmitted to the quarks and leptons by gauge bosons. Since this is a renormalizable theory without fundamental scalars, the naturalness problem is avoided. However, the ETC models that give rise to correct fermion masses have troubles with large FCNC, light pseudo-Goldstone bosons and electroweak precision measurements. Significant attempts to construct realistic ETC models were made recently [10] but phenomenological problems remain to be solved [11, 12].

Although technicolor was introduced as a mechanism for electroweak symmetry breaking which does not depend on the existence of fundamental scalar fields, it is possible to construct technicolor models containing fundamental scalars. Simmons [13] considered a technicolor model with the ordinary fermions receiving mass due to a massive scalar doublet which couples to the technicolor condensate. The phenomenology of this model is acceptable for a large range of parameters [13, 14]. Another possibility is that the scalar doublet is massless while the physical scalar states acquire mass from radiative corrections [15]. The naturalness problem does not appear in technicolor with a scalar doublet

provided the scalar is a composite state, such as a fermion-antifermion state bound by fine-tuned ETC interactions [16]. Samuel [17] considered a supersymmetric[2] version of this model, called bosonic technicolor, which avoids most of the problems of technicolor and of the SSM [17, 19, 20]. As in the case of the SM, these models with scalar doublets offer no insight into the structure of the quark and lepton mass matrices. An exception is a multi-Higgs model [20] with Yukawa couplings controlled by horizontal symmetries.

An explanation for the peculiar pattern of fermion masses might require a mechanism for fermion mass generation not based on the Yukawa couplings of the Higgs doublet. In the mechanism for generating dynamical fermion masses proposed by Kaplan [21] the exchange of technicolored $SU(2)_W$-singlet scalars induces four-fermion effective interactions involving three technifermions and one ordinary fermion. As a result, the ordinary fermions contain an admixture of technibaryon and acquire mass. A hierarchy of masses is produced but the model predicts unacceptable FCNC and tree level contributions to the ρ parameter.

A different attempt to construct a realistic model, in which the exchange of $SU(2)_W$-singlet techniscalars induces four-fermion interactions between two ordinary fermions and two technifermions, is due to Kagan [22]. In this model there are two doublets of technifermions such that the techniscalar exchange contributions to the fermion mass matrices have rank two. Therefore, only two generations are massive at tree level. SUSY is necessary in this model in order to avoid the naturalness problem but also it offers a source of radiative masses for the fermions of the first generation. The hierarchy between the second and third generation should be put in by hand, as in the SM. However, since the fermion masses are quadratic in Yukawa coupling constants, the fine-tuning of the Yukawa couplings is less problematic in the model of Kagan than in the SM. Phenomenological issues associated with $SU(2)_W$-singlet techniscalars and different scenarios for quark mass generation in non-SUSY theories are discussed in Ref. [23].

In this paper we propose a model in which the mass hierarchy between the three generations of quarks and leptons arises naturally.

The model has several features in common with the model of Kagan. There is no Higgs doublet and the $SU(2)_W \times U(1)_Y$ symmetry is broken by technicolor interactions. We introduce technicolored scalar fields which are $SU(2)_W$-singlets, in order to couple the ordinary fermions to the technifermions. When the technifermions condense, the quarks

[2] Earlier attempts of combining SUSY and technicolor can be found in Ref. [18].

and leptons acquire mass. Since there are fundamental scalars in the model, their masses should be protected by SUSY against quadratic divergences.

However, our model is more economic and more natural. There is only one doublet of technifermions. The flavor structure of this supersymmetric technicolor (SUSY-TC) model leads to a realistic pattern of fermion masses. The reason is that the Yukawa couplings of the techniscalars can provide mass only for one generation of fermions while the other two generations acquire smaller masses due to radiative corrections involving gauginos, squarks and sleptons. Such radiative fermion masses were discussed previously [24, 5, 20] but in those cases the "chirality-flip" mixing of squarks or sleptons was produced by Higgs couplings. In our model the interaction of the squarks and sleptons with the technifermions is at the origin of the chirality-flip mixing (a similar mechanism is employed in Ref. [22]). The hierarchy between the second and first generations of fermions is dictated by the structure of the squark and slepton mass matrices, which, in turn, is suggested by the constraints on FCNC. So far, the model has a viable phenomenology with distinctive low energy predictions regarding the fermions of the third generation.

In Section 2 we describe the model. We estimate the fermion masses in Section 3. In Section 4 we discuss the constraints on squark masses from FCNC and we study CP violation effects. The main ideas are summarized in Section 5.

2 The Model

The gauge group is that of a minimal technicolor model: $SU(N_{TC})\times SU(3)_C\times SU(2)_W\times U(1)_Y$. The only source of electroweak symmetry breaking is the vacuum expectation value of the technifermion bilinear. This condensate couples to the weak gauge bosons which become massive. An ordinary fermion has to couple to the condensate in order to acquire mass. This can be done, as we will show in the discussion of quark and lepton masses, by introducing an $SU(2)_W$-singlet scalar which has Yukawa interactions with the ordinary fermion and a technifermion. Such vertices are allowed by Lorentz invariance and gauge symmetry if the left-handed technifermions are $SU(2)_W$-singlets and the right-handed technifermions form doublets. To minimize the radiative electroweak correction parameter S [25] we introduce only one doublet of technifermions.

We consider a low energy theory with global $N = 1$ SUSY broken softly. The Yukawa interactions of the technifermions with the ordinary fermions appear in the superpotential which is expressed only in terms of left-handed chiral superfields. Thus, it is not possible

to have the same scalar involved in the Yukawa interactions of both left-handed and right-handed fermions. However, the scalar superpartners of the left-handed and right-handed components of an $SU(2)_W$-singlet technifermion couple, respectively, to the left-handed and the right-handed fermions; their mixing induces fermion masses. The same pair of techniscalars couples to both the up-type and down-type quarks. The gauge symmetry requires a different pair of techniscalars to couple to the charged leptons.

The charges of the technicolored particles are uniquely determined by imposing hypercharge conservation and the cancellation of the gauge anomalies. The technicolored chiral superfields and their $SU(N_{TC}) \times SU(3)_C \times SU(2)_W \times U(1)_Y$ representations are listed below; the doublet of technifermions responsible for electroweak symmetry breaking is contained in

$$\Upsilon \ : \ (N_{TC}, 1, 2)_0 \ , \qquad \left\{ \begin{array}{l} p_L \ : \ (\bar{N}_{TC}, 1, 1)_{-1} \\ m_L \ : \ (\bar{N}_{TC}, 1, 1)_1 \end{array} \right. \ , \tag{2.1}$$

where $\Upsilon = \left(\begin{array}{c} p^c \\ m^c \end{array} \right)$; the scalar components of the $SU(2)_W$-singlet superfields

$$\begin{array}{rcl} \chi_L \ : \ (\bar{N}_{TC}, 1, 1)_1 & , & \chi^c \ : \ (N_{TC}, 1, 1)_{-1} \ , \\ \phi_L \ : \ (\bar{N}_{TC}, \bar{3}, 1)_{-\frac{1}{3}} & , & \phi^c \ : \ (N_{TC}, 3, 1)_{\frac{1}{3}} \end{array} \tag{2.2}$$

communicate the electroweak symmetry breaking to the leptons and quarks, respectively. The superscript c denotes the charged conjugated superfields. The techni-singlet chiral superfields are those of the MSSM without the Higgs sector:

$$\begin{array}{ll} q_i = \left(\begin{array}{c} u_{i_L} \\ d_{i_L} \end{array} \right) \ , & u_i^c \ , \ d_i^c \ , \\ l_i = \left(\begin{array}{c} \nu_{i_L} \\ e_{i_L} \end{array} \right) \ , & e_i^c \ , \end{array} \tag{2.3}$$

where $i = 1, 2, 3$ is a generation index.

We mention that in a model with one family of technifermions it is enough to introduce only one pair of techni-scalars which are $SU(2)_W$-singlets in order to give masses to both quarks and leptons. The drawback of such a model is that it contains four doublets which produce a large contribution to the S parameter [25]. Note, also, that in the one-doublet technicolor model presented here there are no pseudo-Goldstone bosons. In more complicated models, such as the two-doublet model of Kagan [22], there are pseudo-Goldstone bosons which require additional fields and interactions in order to become massive.

The supersymmetric part of the Lagrangian contains kinetic terms for all the fields, four-scalar interactions proportional to the gauge coupling constants, Yukawa interactions of the fermion and scalar of each chiral supermultiplet with the associated gauginos, and the superpotential. In addition, there are soft SUSY breaking terms [26] consisting of bilinear and trilinear scalar terms and mass terms for the gauginos. Only gauge invariant terms which conserve the baryon number B and the lepton number L are allowed in the Lagrangian. We assign the following B and L numbers to the technicolored chiral superfields:

$$\begin{aligned} p_{{}_L},\, m_{{}_L} \; &: \; L = 0 \,,\; B = 0 \\ \chi_{{}_L} \; &: \; L = -1 \,,\; B = 0 \\ \phi_{{}_L} \; &: \; L = 0 \,,\; B = -1 \end{aligned} \tag{2.4}$$

Apparently, the superpotential includes interactions of all three generations of technisinglet chiral superfields with the technicolored superfields. However, after performing an appropriate unitary transformation in the flavor space, only one generation (the third one, by definition) couples to the technicolored superfields; this is possible because the superpotential is linear in the quark and lepton superfields. Therefore, the most general superpotential is given by

$$\begin{aligned} W \;\; = \;\; & - C_q \epsilon^{\alpha\beta} q_{3_\alpha} \Upsilon_\beta \phi_{{}_L} - C_t u_3^c m_{{}_L} \phi^c + C_b d_3^c p_{{}_L} \phi^c - C_l \epsilon^{\alpha\beta} l_{3_\alpha} \Upsilon_\beta \chi_{{}_L} + C_\tau e_3^c p_{{}_L} \chi^c \\ & + m_\chi \chi_{{}_L} \chi^c + m_\phi \phi_{{}_L} \phi^c + \mathrm{h.c.} \;, \end{aligned} \tag{2.5}$$

where α, β are $SU(2)_W$ indices, $\epsilon^{\alpha\beta}$ is an antisymmetric tensor, m_χ and m_ϕ are mass parameters. The signs in front of the trilinear terms correspond to positive coupling constants $C_f (f = q, t, b, l, \tau)$ in the expressions for fermion masses.

Expressed in terms of scalar and fermion components, the superpotential consist of Yukawa couplings, four-scalar operators, and mass terms for the $SU(2)_W$-singlet technicolored fermions and scalars.

The flavor redefinition we performed to obtain the superpotential given by Eq. (2.5) is an $[SU(3) \times U(1)]^5$ transformation where there is an $SU(3) \times U(1)$ factor for each of the five chiral superfields shown in Eq. (2.3). To find the consequences of such a transformation, it is useful to classify the interactions in terms of a global $[SU(3) \times U(1)]^5$ flavor symmetry. Each of the Yukawa terms in the superpotential breaks one of the $SU(3) \times U(1)$ symmetries down to $SU(2) \times U(1)$. These are the only supersymmetric

interactions that break flavor symmetry. The three-scalar soft terms are linear in squark and slepton fields (see Appendix A) and they also break the flavor symmetry in the scalar sector down to $[SU(2)\times U(1)]^5$. In general, the coefficients of the soft SUSY breaking terms are not related to the Yukawa coupling constants so that the combination of three-scalar terms and Yukawa interactions breaks the flavor symmetry down to $[U(1)]^5$. The squark and slepton mass terms break completely the flavor symmetry in the scalar sector. Therefore, the flavor transformation changes only the coefficients of the soft SUSY breaking terms which involve scalars. In particular, the squark mass matrices are redefined by unitary transformations.

3 Fermion Mass Generation

We begin the discussion of fermion mass generation by studying the interactions responsible for the top mass. We use the same notation for the fermions as the one in Eqs. (2.1)-(2.3) for the corresponding chiral superfields and we switch from two-component to four component spinors. We denote the scalars by the symbols used for their fermion partners with a tilde and we define "right-handed" scalar fields: $\tilde{\phi}_R \equiv \tilde{\phi}^{c\dagger}$, etc. In the case of the third generation fermions, we use the conventional notation, t, b, τ, without distinguishing between weak and mass eigenstates.

The top quark has Yukawa interactions with the m technifermion:

$$C_q \bar{m}_R t_L \tilde{\phi}_L + C_t \bar{t}_R m_L \tilde{\phi}_R^\dagger + \text{h.c.} \tag{3.1}$$

Both the soft SUSY breaking terms and the mass terms in the superpotential contribute to the $\tilde{\phi}_L$-$\tilde{\phi}_R$ mass matrix (see Eq. (A.2)). At energies lower than the masses of $\tilde{\phi}_L$ and $\tilde{\phi}_R$ the exchange of these scalars gives rise to four-fermion operators involving two quarks and two technifermions. The top mass arises due to a four-fermion operator which requires $\tilde{\phi}_L$-$\tilde{\phi}_R$ mixing in the techni-scalar exchange, as shown in Fig. 1:

$$\frac{C_q C_t}{\mathcal{M}_\phi^2} \left(\bar{m}_L t_R\right)\left(\bar{t}_L m_R\right) + \text{h.c.}\ , \tag{3.2}$$

where $\mathcal{M}_\phi^2$ is a combination of the diagonal, $(M_{\tilde{\phi}}^2)_{LL}$, $(M_{\tilde{\phi}}^2)_{RR}$, and off-diagonal, $(M_{\tilde{\phi}}^2)_{LR}$, elements of the mass matrix for $\tilde{\phi}_L$ and $\tilde{\phi}_R$,

$$\mathcal{M}_\phi^2 = \frac{(M_{\tilde{\phi}}^2)_{LL}(M_{\tilde{\phi}}^2)_{RR} - (M_{\tilde{\phi}}^2)_{LR}^2}{(M_{\tilde{\phi}}^2)_{LR}}\ . \tag{3.3}$$

Applying a Fierz transformation to the operator (3.2) we find the top mass:

$$m_t = \frac{C_q C_t}{4\mathcal{M}_\phi^2}\langle \bar{m}m \rangle \ . \tag{3.4}$$

According to naive dimensional analysis [27], the condensate is related to the technipion decay constant v by $\langle \bar{m}m \rangle \approx 4\pi v^3$. In one-doublet technicolor models, $v \approx 246\,\text{GeV}$. The constraint on $\mathcal{M}_\phi^2$ imposed by Eq. (3.4) is

$$\mathcal{M}_\phi \approx 0.5\,\text{TeV} \times (C_q C_t)^{\frac{1}{2}} \ , \tag{3.5}$$

where we used $m_t = 176\,\text{GeV}$ [28]. Since perturbation theory requires Yukawa couplings smaller than $\sim 4\pi$, this equation places an upper bound on the masses of the $\tilde{\phi}$ scalars in the absence of fine-tuning in the mass matrix for $\tilde{\phi}_L$ and $\tilde{\phi}_R$. We will assume $\mathcal{M}_\phi \sim 1\text{TeV}$. The technicolor corrections to the diagram shown in Fig. 1 are not important since the constituent masses of the $\tilde{\phi}$ scalars are of the order of the current mass $\mathcal{M}_\phi$. Note that the Yukawa interactions may be treated as a small perturbation on the technicolor dynamics [11] because $m_t/4\pi v$ is small.

The b quark acquires mass similarly, by coupling to the p technifermion. Since the custodial $SU(2)_R$ symmetry requires $\langle \bar{m}m \rangle \approx \langle \bar{p}p \rangle$, the only source for the large mass ratio of the t and b quarks is the ratio of the Yukawa couplings,

$$\frac{m_t}{m_b} = \frac{C_t}{C_b} \ . \tag{3.6}$$

The τ mass is produced by the exchange of $\tilde{\chi}$ scalars, so the t to τ mass ratio depends on the scalar masses:

$$\frac{m_t}{m_\tau} = \frac{C_q C_t}{C_l C_\tau}\left(\frac{\mathcal{M}_\chi}{\mathcal{M}_\phi}\right)^2 \ , \tag{3.7}$$

where $\mathcal{M}_\chi^2$ is a combination of the elements of the mass matrix for $\tilde{\chi}_L$ and $\tilde{\chi}_R$ analogous to Eq. (3.3).

Eq. (3.5) provides some information about the SUSY breaking scale $\mathcal{M}_s$. As long as we do not refer to the high energy theory responsible for SUSY breaking there are no theoretical constraints on the coefficients of the soft SUSY breaking terms. However, we assume these coefficients to have the same order of magnitude, given by $\mathcal{M}_s$. The masses of the $\tilde{\phi}$ scalars indicate

$$\mathcal{M}_s \sim \mathcal{O}(1\,\text{TeV}) \ . \tag{3.8}$$

Note that in general the technifermion masses m_ϕ and m_χ, the techniscalar mixings $(M^2_{\tilde{\phi}})_{LR}$ and $(M^2_{\tilde{\chi}})_{LR}$ and the Yukawa coupling constants $C_f (f = q, t, b, l, \tau)$ are complex

numbers. However, their phases can be absorbed in the scalar and fermion fields so that all the quantities which appear in Eqs. (3.3)-(3.7) are real. The phase redefinition can be done in several ways and introduces new complex phases in other coupling constants. In the discussion of CP violation (see Section 4.2) we will use an explicit phase convention.

Although the quarks of the first and second generations do not couple to the technifermions, there are contributions to their masses from interactions with gauginos and squarks. The electroweak symmetry breaking enters in these radiative masses through the mixing of the left-handed and right-handed squarks. This chirality-flip mixing is produced by the exchange of the ϕ technifermion, whose mass, m_ϕ, is an arbitrary parameter in the superpotential. For simplicity, we will assume $m_\phi \sim \mathcal{M}_s$.

In the "super-weak" basis, where the quark-squark-gluino vertices are flavor diagonal and the superpotential is given by Eq. (2.5), only the squarks of the third generation couple to the technifermions. However, in this basis the left-handed and right-handed squark mass matrices are in general non-diagonal. The off-diagonal squark masses combined with the $\tilde{u}_{3_L}$-$\tilde{u}_{3_R}$ mixing produce chirality-flip mixings of the $\tilde{u}_1$ and $\tilde{u}_2$ squarks. The quark-squark-gluino interaction leads to the one-loop graph shown in Fig. 2 which yield an effective four-fermion interaction. When the technifermions condense, this graph makes the largest contribution to the elements of the up quark mass matrix:

$$m^u_{ij} = \frac{\alpha_s}{\pi} m_t \frac{m_\phi \mathcal{M}_\phi^2}{m_{\tilde{g}}^3} f^u_{ij} \ , \tag{3.9}$$

where $i, j = 1, 2, 3$, f^u_{ij} are functions of the squark and gluino masses given in Appendix B, $\alpha_s \approx 0.1$ is the strong coupling constant at a scale $\sim \mathcal{M}_s$ and $m_{\tilde{g}}$ is the gluino mass. A rough estimate (see Eq. (B.2)) gives $|f^u_{ij}| \lesssim 10^{-1}$. Hence, the m^u_{33} element of the quark mass matrix is given by Eq. (3.4) ($m^u_{33} \approx m_t$) while the other elements are much smaller,

$$\frac{|m^u_{ij}|}{m_t} \leq \mathcal{O}(10^{-2}) \ . \tag{3.10}$$

One of the benefits of this structure of the quark mass matrix is that analytical expressions for the quark masses may be obtained. Diagonalizing the mass matrix we find the quark masses to first order in m^u_{ij}/m_t:

$$\begin{aligned} m_c &= \left(|m^u_{11}|^2 + |m^u_{12}|^2 + |m^u_{21}|^2 + |m^u_{22}|^2\right)^{\frac{1}{2}} \\ m_u &= \frac{|m^u_{11} m^u_{22} - m^u_{12} m^u_{21}|}{(|m^u_{11}|^2 + |m^u_{12}|^2 + |m^u_{21}|^2 + |m^u_{22}|^2)^{\frac{1}{2}}} \ . \end{aligned} \tag{3.11}$$

Since the supersymmetric part of the Lagrangian has a flavor $[SU(2)\times U(1)]^5$ symmetry with respect to the first and second generations, the super-weak basis is defined up to such a transformation. Therefore, there is a super-weak basis in which the $(1,2)$ elements of the left-handed and right-handed squark mass matrices, M_L^{q2} and M_R^{u2}, vanish. The complex phases of the other non-diagonal elements can be absorbed in the definition of the squark fields, such that the squark mass matrices are real and symmetric. We will assume the $(2,3)$ elements and the diagonal elements to be of order $\mathcal{M}_s^2$. As we will discuss in Section 4.1, the constraints from FCNC require small values of the $(1,3)$ elements:

$$
\begin{aligned}
\epsilon_L^q &\equiv \frac{(M_L^{q2})_{13}}{(M_L^{q2})_{23}} \lesssim \mathcal{O}(10^{-1}) \\
\epsilon_R^u &\equiv \frac{(M_R^{u2})_{13}}{(M_R^{u2})_{23}} \lesssim \mathcal{O}(10^{-1})
\end{aligned}
\tag{3.12}
$$

This structure of the squark mass matrices allow us to obtain the unitary matrices U_L^q and U_R^u which transform the squark fields from the super-weak basis to the mass eigenstate basis; to first order in ϵ_L^q,

$$
U_L^q = \begin{pmatrix} 1 & \mathcal{O}(\epsilon_L^q) & \mathcal{O}(\epsilon_L^q) \\ \mathcal{O}(\epsilon_L^q) & \cos\theta_L & -\sin\theta_L \\ \mathcal{O}(\epsilon_L^q) & \sin\theta_L & \cos\theta_L \end{pmatrix} , \tag{3.13}
$$

where

$$
\cos^2\theta_L = \frac{1}{2}\left\{1+\left|(M_L^{q2})_{22}-(M_L^{q2})_{33}\right|\left[\left((M_L^{q2})_{22}-(M_L^{q2})_{33}\right)^2+4(M_L^{q2})_{23}^2\right]^{-1/2}\right\} \tag{3.14}
$$

and similar relations hold for U_R^u. Eqs. (3.13), (B.2) and (3.9) yield the following structure for the up quark mass matrix:

$$
m^u = m_t \begin{pmatrix} \mathcal{O}(\epsilon_L^q\epsilon_R^u\beta^u) & \mathcal{O}(\epsilon_L^q\beta^u) & \mathcal{O}(\epsilon_L^q\beta^u) \\ \mathcal{O}(\epsilon_R^u\beta^u) & \beta^u & \mathcal{O}(\beta^u) \\ \mathcal{O}(\epsilon_R^u\beta^u) & \mathcal{O}(\beta^u) & 1 \end{pmatrix} , \tag{3.15}
$$

where

$$
\beta^u = \frac{\alpha_s}{\pi}\frac{m_\phi \mathcal{M}_\phi^2}{m_{\tilde{g}}^3} f_{22}^u . \tag{3.16}
$$

The up and charm masses can be estimated by combining Eqs. (3.11) and (3.15):

$$
\frac{m_c}{m_t} \sim \beta^u \tag{3.17}
$$

$$
\frac{m_u}{m_c} \sim \mathcal{O}(\epsilon_L^q\epsilon_R^u) \tag{3.18}
$$

These are realistic predictions [29], provided $\epsilon^q_L \epsilon^u_R \sim 10^{-2}$ (see Eq. (3.12)) and $\beta^u \sim 10^{-2}$. Note that the quark masses computed here are at a scale of order $\mathcal{M}_s$ and are smaller by a factor ~ 2 than the masses at a scale of 1 GeV [20]. The quark mass matrix m^u is diagonalized by unitary matrices:

$$V^{u\dagger}_L m^u V^u_R = \mathrm{diag}(m_u, m_c, m_t) \tag{3.19}$$

where, to first order in β^u and ϵ^q_L,

$$V^u_L = \begin{pmatrix} 1 & \mathcal{O}(\epsilon^q_L) & \mathcal{O}(\epsilon^q_L \beta^u) \\ \mathcal{O}(\epsilon^q_L) & 1 & \mathcal{O}(\beta^u) \\ \mathcal{O}(\epsilon^q_L \beta^u) & \mathcal{O}(\beta^u) & 1 \end{pmatrix} ; \tag{3.20}$$

V^u_R has the same structure, with ϵ^u_R instead of ϵ^q_L.

The elements of the down quark mass matrix are given by Eq. (3.9) with m_t replaced by m_b and with different functions f^d_{ij}. The comparatively large strange to bottom mass ratio, β^d, given by Eq. (3.16) with f^u_{22} replaced by f^d_{22}, requires the ratio $M_{\tilde{\phi}}/m_{\tilde{g}}$ to have a rather large value ~ 3 and f^d_{22} to be close to its upper bound $\sim 10^{-1}$. The difference between the m_s/m_b and m_c/m_t ratios is due to the different squark mass matrices which contribute to f^u_{22}, respectively f^d_{22}.

The FCNC constraints on the down squark sector are stronger (see Section 4.1), giving an upper bound

$$\epsilon^q_L \epsilon^d_R \lesssim \mathcal{O}(10^{-3}) , \tag{3.21}$$

where

$$\epsilon^d_R \equiv \frac{(M^{d2}_R)_{13}}{(M^{d2}_R)_{23}} . \tag{3.22}$$

This makes the contribution of the one-loop graph of Fig. 2 (with up-type quarks and squarks replaced by down-type ones) to the down mass, m_d, very small. However, there are other contributions from two-loop graphs in which the $\tilde{d}_{1L}$-$\tilde{d}_{1R}$ mixing is mediated by three-scalar interactions and techni-gluino exchange (see Fig. 3). The coefficient of the $\tilde{d}_{1L}$-$\tilde{d}_{1R}$ mass term produced is of order $\mu_{d_1}\mu_{q_1}/(4\pi)^2$ where μ_{d_1} and μ_{q_1} are the mass coefficients of the $\tilde{d}\tilde{\phi}\tilde{p}$ soft SUSY breaking terms (see Eq. (A.5)). The down quark mass produced is large enough provided $\mu_{d_1}\mu_{q_1} \sim \mathcal{M}_s^2$.

The down quark mass matrix is diagonalized by unitary matrices, V^d_L and V^d_R, with the same structure as V^u_L. Therefore, the CKM matrix

$$V_{KM} = V^{u\dagger}_L V^d_L \tag{3.23}$$

has also the structure shown in Eq. (3.20), with elements $V_{us} \sim \mathcal{O}(10^{-1})$, $V_{cb} \sim \mathcal{O}(10^{-2})$ and $V_{cb} \sim \mathcal{O}(10^{-3})$.

In the case of charged leptons, the elements of the mass matrix are given by one-loop graphs similar to the one in Fig. 2 with the gluino, the ϕ technifermion and the squarks replaced, respectively, by a zino or photino, a χ technifermion and sleptons. The m_μ/m_τ ratio differ from m_c/m_t by a factor $\sim (\alpha_2/\alpha_s)(\mathcal{M}_\chi^2/\mathcal{M}_\phi^2)$ where α_2 is the weak coupling constant at a scale $\sim \mathcal{M}_s$. Since $m_t/m_\tau \approx 100$, Eq. (3.7) indicates a large $\mathcal{M}_\chi^2/\mathcal{M}_\phi^2$ ratio; thus, the ratio $m_\mu/m_\tau \approx \frac{1}{15}$ can be readily obtained. If the charged slepton mass matrices and the squark mass matrices have a similar structure, as it is suggested by the constraints from $\mu \to e\gamma$ [4], then the electron mass is predicted to be two orders of magnitude smaller than the muon mass. The neutrinos remain massless because we did not introduce right-handed spinors.

In conclusion, the mass hierarchy between the fermion generations is established. It is remarkable that the SUSY-TC model is able to reproduce the complicated pattern of fermion masses with only few assumptions about the soft SUSY breaking terms and the parameters in the superpotential.

4 Flavor-Changing Neutral Currents

The measurements of FCNC effects impose severe constraints on ETC models and on SUSY models. Therefore, FCNC represent an important test for a SUSY-TC model. In this section we discuss the FCNC in our model, concentrating on the quark sector.

4.1 Neutral meson mixing

As we showed in Section 2, in the super-weak basis only the b and t quarks couple to the technifermions. However, quark mixings are produced at the one loop level and, as a result, the quarks of the first and second generations in the mass eigenstate basis have Yukawa interactions with the technicolored fields proportional to the small mixing angles of the third generation. Thus, box diagrams with techniscalars and technifermions in the internal lines contribute to $K - \bar{K}$ and $B - \bar{B}$ mixing. Nevertheless, these contributions are suppressed by a factor of order $\mathcal{M}_s^2/M_W^2$ with respect to the SM amplitudes and can be ignored. Other contributions to the $\Delta S = 2$ or $\Delta B = 2$ amplitudes involving technicolored fields are given by dimension-12 operators and are much smaller.

Larger FCNC are produced due to the techni-singlet sparticles. In generic SUSY models [4], the quark and squark mass matrices are diagonalized by different transformations. Therefore, the quark-squark-gaugino vertices are flavor non-diagonal in the mass eigenstate basis and give rise to FCNC. It is convenient to compute the FCNC effects in the super-KM basis [30], where the quark-squark-gluino vertices are flavor diagonal, using the mass insertion approximation. This procedure was used extensively [4] to put bounds on the off-diagonal elements of the left-handed and right-handed squark mass matrices and on the chirality-flip mixings of squarks belonging to different generations. The tightest bounds are on down-type squark mixings and come from gluino one-loop diagrams contributing to $b \to s\gamma$ and to $K - \bar{K}$ and $B - \bar{B}$ mixing.

In our model, the chirality-flip mixing of the down-type squarks arises due to the ϕ technifermion exchange diagram shown in Fig. 4. At energies below m_ϕ, the effect of ϕ exchange may be approximated by local operators of dimension-five. The rules of naive dimensional analysis show that this is a good approximation if

$$m_\phi > v\frac{C_q C_b}{4\pi} \ . \tag{4.1}$$

When the technifermions condense, the mixing of the scalar-bottoms in the super-weak basis is given by

$$(M_{LR}^{d\,2})_{33} = -\frac{C_q C_b}{2m_\phi}\langle \bar{p}p\rangle \ . \tag{4.2}$$

Combining Eqs. (3.4), (3.6) and (4.2) gives

$$\left|(M_{LR}^{d\,2})_{33}\right| = 2m_b\frac{\mathcal{M}_\phi^2}{m_\phi} \ , \tag{4.3}$$

which is a small mixing: $(M_{LR}^{d\,2})_{33}/\mathcal{M}_s^2 \sim 10^{-2}$. The chirality-flip mixings of the down-type squarks belonging to different generations are composed of a $(3,i)$ or $(i,3)$ $(i = 1,2)$ element of $M_L^{q\,2}$ or $M_R^{d\,2}$ and the $\tilde{d}_{3_L}$-$\tilde{d}_{3_R}$ mixing given by Eq. (4.3). The chirality-flip mixing produced by techni-gluino exchange and trilinear scalar interactions, as in Fig. 3, is also small:

$$\frac{\left|(M_{LR}^{d\,2})_{ij}\right|}{\mathcal{M}_s^2} \sim \mathcal{O}(10^{-2}) \ , \tag{4.4}$$

where $i,j = 1,2,3$. In the super-KM basis, the $\tilde{d}_1$ and $\tilde{d}_2$ squarks couple to the technifermions with Yukawa couplings suppressed by the b-d and b-s quark mixing angles, respectively. This also produces very small chirality-flip squark mixings. We conclude that the stringent bounds from $K - \bar{K}$, $B - \bar{B}$ mixing and $b \to s\gamma$ on the chirality-flip squark mixing [4] are naturally satisfied in our model.

There are, however, important constraints on the $(M_L^{q2})_{ij}$ and $(M_R^{d2})_{ij}$ $(i \neq j)$ mixings due to the gluino box diagrams contributing to the $K - \bar{K}$ mass difference [4]

$$\frac{\left|(M_L^{d2})'_{12}\right|}{\mathcal{M}_s^2}, \; \frac{\left|(M_R^{d2})'_{12}\right|}{\mathcal{M}_s^2} \leq \mathcal{O}(10^{-1}) \frac{\mathcal{M}_s}{1\mathrm{TeV}}$$
$$\left| \frac{(M_L^{d2})'_{12}}{\mathcal{M}_s^2} \frac{(M_R^{d2})'_{12}}{\mathcal{M}_s^2} \right|^{1/2} \leq \mathcal{O}(10^{-2}) \frac{\mathcal{M}_s}{1\mathrm{TeV}} \tag{4.5}$$

and to the $B - \bar{B}$ mass difference

$$\frac{\left|(M_L^{d2})'_{13}\right|}{\mathcal{M}_s^2}, \; \frac{\left|(M_R^{d2})'_{13}\right|}{\mathcal{M}_s^2} \leq \mathcal{O}(10^{-1}) \frac{\mathcal{M}_s}{1\mathrm{TeV}} \tag{4.6}$$

where the primed matrix elements refer to the super-KM basis.

The constraints on chirality-conserving mixing between the squarks of the second and third generations are loose. The $B_s - \bar{B}_s$ mass difference is expected to be larger by an order of magnitude than the $B - \bar{B}$ mass difference [31]. Combined with Eq. (4.6), this leads to

$$\frac{\left|(M_L^{d2})'_{23}\right|}{\mathcal{M}_s^2}, \; \frac{\left|(M_R^{d2})'_{23}\right|}{\mathcal{M}_s^2} \leq \mathcal{O}(1) \frac{\mathcal{M}_s}{1\mathrm{TeV}} \; . \tag{4.7}$$

The up squark mass matrices are also constrained; the gluino box diagrams contributing to the $D - \bar{D}$ mass difference [4] require

$$\frac{\left|(M_L^{u2})'_{12}\right|}{\mathcal{M}_s^2}, \; \frac{\left|(M_R^{u2})'_{12}\right|}{\mathcal{M}_s^2} \leq \mathcal{O}(10^{-1}) \frac{\mathcal{M}_s}{1\mathrm{TeV}} \; . \tag{4.8}$$

The wino box diagrams contributing to the $K - \bar{K}$ mass difference give an upper limit for the squark mixing which is larger than the bound in Eq. (4.8).

The small values of the ratios in Eqs. (4.5), (4.6) and (4.8) are unnatural. In general one expects these ratios to be of order one [5]. A possible solution to this problem might be the existence of gauged horizontal symmetries [7]. Note that the bounds given by Eqs. (4.5)-(4.8) are $\sim \mathcal{M}_s$ which implies looser bounds in our model than in the SSM where the SUSY breaking scale is likely to be below 1 TeV.

The $(1, 2)$ elements of the down squark mass matrices in the super-KM basis are related to the ones in the super-weak basis by:

$$\left|(M_L^{d2})'_{12}\right| \approx \epsilon_L^q \left|(M_L^{q2})_{11} - (M_L^{q2})_{22}\right|$$
$$\left|(M_R^{d2})'_{12}\right| \approx \epsilon_R^d \left|(M_R^{d2})_{11} - (M_R^{d2})_{22}\right| \tag{4.9}$$

where we kept only the leading terms in β^d, ϵ_L^q and ϵ_R^d. The relation between the Cabbibo angle and $(M_L^{q2})_{12}$, given by Eq. (3.20), requires $\epsilon_L^q \sim \mathcal{O}(10^{-1})$. Therefore, ϵ_R^d is strongly constrained by Eq. (4.5). In order to avoid excessive fine-tuning of $(M_R^{d2})_{13}$, we assume a reasonably small value of

$$\left[\frac{(M_L^{q2})_{11} - (M_L^{q2})_{22}}{\mathcal{M}_s^2}\right]\left[\frac{(M_R^{d2})_{11} - (M_R^{d2})_{22}}{\mathcal{M}_s^2}\right] \sim \mathcal{O}(10^{-1}) \ . \tag{4.10}$$

In this case, the bound in Eq. (4.5) is saturated if $\epsilon_R^d \sim \mathcal{O}(10^{-2})$. This is also a sufficient condition for satisfying the constraints on the $(1,3)$ elements given by Eq. (4.6).

In the up squark sector, the relation between the super-KM and the super-weak basis is given by Eq. (4.9), with the upper index d replaced by u. The inequalities (4.8) are satisfied if $\epsilon_R^u, \epsilon_L^q \lesssim \mathcal{O}(10^{-1})$.

It is interesting that the bounds on the $(1,2)$ elements of the squark mass matrices in the super-KM basis correspond to bounds on the $(1,3)$ elements in the super-weak basis.

4.2 CP violation

To keep track of the relative complex phases of the coupling constants relevant for CP violation we will adopt the phase convention described below. The masses m_ϕ and m_χ become real by absorbing their phases in the fermion fields ϕ_L and χ_L. The phases of the off-diagonal mass terms $(M^2_{\tilde{\phi}})_{LR}$ and $(M^2_{\tilde{\chi}})_{LR}$ are absorbed in the scalar fields $\tilde{\chi}_L$ and $\tilde{\phi}_L$. These redefinitions introduce new phases in the ϕ_L-$\tilde{\phi}_L$-gaugino and χ_L-$\tilde{\chi}_L$-gaugino interactions and also in the Yukawa couplings from the superpotential. The phases of the C_q Yukawa coupling from the quark and squark sectors are now different and they are absorbed in the fermion Υ, respectively scalar $\tilde{\Upsilon}$ doublets, while a new phase appears in the Υ-$\tilde{\Upsilon}$-gaugino interactions. A redefinition of the u_3^c, d_3^c and e_3^c superfields yields C_t, C_b and C_τ real. The C_l coupling constant has also different phases in the lepton and slepton vertices; these are absorbed by the l_3 and $\tilde{l}_3$ fields leading to a phase in the l_3-$\tilde{l}_3$-gaugino interactions.

At this stage, the only complex coupling constants left are in the soft SUSY breaking terms and in the fermion-scalar-gaugino interactions mentioned above. The gaugino masses are in general complex. When the gaugino field is redefined to have real mass, a new phase, $\delta_{\tilde{g}}$, is introduced in the quark-squark-gaugino vertices. If an internal gaugino line connects quarks of the same chirality, the phases introduced in the two vertices cancel each other. Nevertheless, the complex phases of the gaugino masses are relevant when

the gaugino propagator connects quarks of different chiralities, as it is the case in the diagram shown in Fig. 2. Thus, there are contributions linear in $\delta_{\tilde{g}}$ to the neutron dipole moment (NDM) from the one-loop diagram shown in Fig. 2 with an external photon line attached to one of the internal lines. These are similar with the SSM contributions to the NDM [32]. The experimental limit on the NDM [29] requires $\delta_{\tilde{g}} \lesssim 10^{-2}$. Also, there are corrections of order $\delta_{\tilde{g}}$ to the phases of all the elements of a quark mass matrix except the (3,3) one. However, the CKM elements can be expressed in terms of quark mass ratios such that the phase $\delta_{\tilde{g}}$ is largely canceled out. Therefore, we will ignore $\delta_{\tilde{g}}$ in the discussion of the phases in the quark mass matrices.

A squark mass matrix is hermitian and has three complex phases. However, in the super-weak basis, the $(1,2)$ elements of the squark mass matrices vanish and, therefore, there are only two phases left. These can be absorbed in the squark fields of the first two generations. The coupling constants of the quark-squark-gaugino vertices are kept real by including the same phases in the definition of the corresponding quark fields.

The result of the above phase convention is that there is no contribution from the squark mixings to the phase of the diagram shown in Fig. 2. Thus, the leading contributions to the quark mass matrices (see Eqs. (3.4) and (3.9)) are real.

This result has interesting consequences. The origin of CP violation should be rather different than the one in the SM since the CKM matrix is approximately real. Also, real quark mass matrices are relevant for the strong CP problem. The strong CP parameter $\bar{\theta}$ receives in this case no contribution from the quark mass matrix [1] while the QCD contributions can be small enough if CP is spontaneously broken. However, the complex phases of the quark mass matrices that we describe below give corrections to $\bar{\theta}$ much larger than the experimental limit of 10^{-9}, such that the strong CP problem persists.

Complex phases in the quark mass matrices come from additional loops involving trilinear scalar interactions or gaugino-technifermion-techniscalar interactions. We will denote generically the complex phases of the coupling constants of these interactions by Δ. The two-loop corrections to the (2,2), (2,3), (3,2) and (3,3) elements of the quark mass matrices are small; the typical size of the phases of these elements is $\sim 10^{-2}\Delta$. The other elements are smaller (see Eq. (3.20)), due to the structure of the squark mass matrices. Thus, the two-loop contributions to the imaginary parts of these elements are larger; writing $\mathrm{Arg}\left(m_{ij}^{u,d}\right) = \delta_{ij}^{u,d}$, we estimate:

$$\delta_{11}^{u}, \delta_{11}^{d}, \delta_{12}^{d}, \delta_{21}^{d}, \delta_{13}^{d}, \delta_{31}^{d} \sim \Delta$$

$$\delta^u_{12}, \delta^u_{21}, \delta^u_{13}, \delta^u_{31} \sim 10^{-1}\Delta \tag{4.11}$$

To see the effect of these phases, we consider the Wolfenstein parametrization of the CKM matrix [33], keeping only the first non-vanishing terms of the expansion in $\lambda = \sin\theta_c$, where θ_c is the Cabbibo angle:

$$V_{KM} = \begin{pmatrix} 1 & \lambda & A\lambda^3(\rho - i\eta) \\ -\lambda & 1 & A\lambda^2 \\ A\lambda^3(1-\rho-i\eta) & -A\lambda^2 & 1 \end{pmatrix} . \tag{4.12}$$

Computing V_{KM}, as discussed in Section 3, to first order in $\beta^u, \beta^d, \epsilon^q_L$, gives:

$$\begin{aligned} \lambda &= \left|\frac{m^d_{12}}{m^d_{22}} - \frac{m^u_{12}}{m^u_{22}}\right| \sim \mathcal{O}(\epsilon^q_L) \\ A\lambda^2 &= \left|\frac{m^d_{23}}{m^d_{33}} - \frac{m^u_{23}}{m^u_{33}}\right| \sim \mathcal{O}(\beta^d) \end{aligned} \tag{4.13}$$

and more complicated expressions for η and ρ, which, together with the estimated phases of the quark mass terms, indicate

$$\begin{aligned} \eta &\sim \mathcal{O}(10^{-1}\Delta) \\ |\rho(1-\rho)| &\sim \mathcal{O}(1) \end{aligned} \tag{4.14}$$

The measurements of CP asymmetry in semileptonic decays of K_L show [34]

$$\frac{\mathrm{Im}M_{12}}{\mathrm{Re}M_{12}} \approx 6.5 \times 10^{-3} , \tag{4.15}$$

where M_{12} is the off-diagonal element of the $K - \bar{K}$ mass matrix. If the phase from the CKM matrix is solely responsible for CP violation, as it is in the SM, then Eq. (4.15) requires [31]

$$\begin{aligned} 0.2 \lesssim \eta_{SM} &\lesssim 0.6 \\ |\rho_{SM}| &\lesssim 0.4 \end{aligned} \tag{4.16}$$

Comparing Eqs. (4.14) and (4.16), we conclude that the CP violation provided by the CKM matrix is not sufficient in our model; hence, the bulk of CP violation in $K - \bar{K}$ mixing is due to SUSY-box diagrams. The relevant phases are those of the $(M^{d^2}_L)'_{12}$ and $(M^{d^2}_R)'_{12}$ squark mixing in the super-KM basis. We will denote generically these squark mixings and their phases by M^2_{ds} and Δ_{ds} respectively. These appear in the gluino box diagrams:

$$\mathrm{Im}M_{12} \approx \tan(2\Delta_{ds})\mathrm{Re}(M_{12})_{gluino} . \tag{4.17}$$

Eq. (4.5) can be written as

$$\frac{\mathrm{Re}(M_{12})_{gluino}}{\mathrm{Re}M_{12}} \sim \left(10^2 \frac{|M^2_{ds}|}{\mathcal{M}^2_s}\right)^2 \ , \tag{4.18}$$

which combined with Eqs. (4.15) and (4.17) gives

$$\left(10^2 \frac{|M^2_{ds}|}{\mathcal{M}^2_s}\right)^2 \tan(2\Delta_{ds}) \sim \mathcal{O}(10^{-2}) \ . \tag{4.19}$$

If the mass ratios in Eq. (4.5) are close to their limits, then Eq. (4.19) indicates the size of the phases of the squark mass mixings:

$$\Delta_{ds} \sim \mathcal{O}(10^{-2}) \ . \tag{4.20}$$

An explicit computation of the squark mass matrices in the super-KM basis, involving unitary transformations with the V^d_L and V^d_R matrices on the squark mass matrices in the super-weak basis, shows $\Delta_{ds} \sim \delta^d_{12}$. Comparing, then, Eqs. (4.11) and (4.20), we obtain the size of the phases of the three-scalar interactions:

$$\Delta \sim \mathcal{O}(10^{-2}) \ . \tag{4.21}$$

Such small phases might arise naturally if there is spontaneous CP violation [35].

The SM predicts large CP asymmetries in B meson decays [36] because the phase in the CKM matrix is $\mathcal{O}(1)$ (see Eq.(4.16)). Eqs. (4.14) and (4.21) show that the situation is totally different in our model: the phases responsible for CP violation in B meson decays are of order

$$\eta \sim \mathcal{O}(10^{-3}) \ . \tag{4.22}$$

Note that the one-loop diagrams with sparticles in internal lines give small contributions to the B decays. In particular, the $B-\bar{B}$ mixing amplitude given by SUSY box diagrams is small because of the severe bounds on ϵ^q_L and ϵ^d_R from $K-\bar{K}$ mixing (see Eqs. (4.5), (4.6) and (3.21)). Hence, the mechanism for CP violation in our model is the same as in the SM, but the effects are smaller by a factor $\eta_{SM}/\eta \sim 10^2$. For example, we estimate the size of CP asymmetry in $B \to \Psi K_S$ [36] to be

$$-\,\mathrm{Im}\left[\frac{q}{p}\frac{A(\bar{B}^0 \to \Psi K_S)}{A\,(B^0 \to \Psi K_S)}\right] \approx \frac{2\eta}{1-\rho} \sim \mathcal{O}(10^{-3}) \ . \tag{4.23}$$

where q/p is the $B-\bar{B}$ mixing parameter. This value is two orders of magnitude smaller than the resolution of the proposed experiments [37] on CP violation in B decays.

Another consequence of Eq. (4.22) is a small direct CP violation in $K_L \to \pi\pi$ decays. To see this, note that in the SM the ε'/ε parameter is proportional to η_{SM} [38]. Therefore, a scaling of the SM result gives $\varepsilon'/\varepsilon \sim \mathcal{O}(10^{-6})$ in the SUSY-TC model. The contributions from SUSY penguin diagrams do not exceed this value. Such a small value is inconsistent with the result of the CERN NA31 experiment ($\varepsilon'/\varepsilon = (23 \pm 7) \times 10^{-4}$) [39] but is consistent with the result of the Fermilab E731 experiment ($\varepsilon'/\varepsilon = (7.4 \pm 6.0) \times 10^{-4}$) [40].

5 Conclusions

We have proposed a supersymmetric one-doublet technicolor model with the superpotential containing Yukawa couplings of the quarks and leptons of the third generation with a technifermion and a $SU(2)_W$-singlet techniscalar. These interactions give rise to masses for the fermions of the third generation. The fermions of the other generations have radiative masses such that a correct mass hierarchy arises. However, the model offer no insight into the origin of the large top to bottom mass ratio, given by a ratio of Yukawa coupling constants. In order to obtain a realistic top mass, the SUSY breaking scale should be ~ 1 TeV. In the low energy SUSY theory, the sparticle masses are not determined; consequently, it is not possible to make more precise predictions for the fermion mass matrices.

The contributions of the technicolored particles to $K-\bar{K}$ and $B-\bar{B}$ mixing are small. Comparing with the SSM, the amount of fine-tuning in the squark mass matrices required to avoid large FCNC is slightly reduced in our model.

With an appropriate redefinition, the complex phase in the CKM matrix is $\mathcal{O}(10^{-3})$. The main contributions to CP violation in $K-\bar{K}$ mixing come from gluino box diagrams. The mechanisms for CP violation in B meson decays and for direct CP violation in K_L decays are the same as in the SM. The CP asymmetries in B decays are smaller by two orders of magnitude than the asymmetries predicted in the SM and will not be detected at the proposed B-factories. Also, the CP asymmetry in $\Delta S = 1$ transitions is tiny.

To decide whether the model is viable it is necessary to explore many other phenomenological issues: electroweak precision measurements, FCNC in the lepton sector, constraints from cosmology, etc. Also, it is interesting to study how this SUSY-TC model fits into a high energy theory, such as grand unification or supergravity.

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Appendix A

This Appendix presents the soft SUSY breaking terms. We use the notation described at the beginning of Section 3.

The techniscalar mass terms can be written

$$\mathcal{L}_{2s} = M^2_{\tilde{\Upsilon}} \tilde{\Upsilon}^\dagger \tilde{\Upsilon} + M^2_{\tilde{p}_L} \tilde{p}_L^\dagger \tilde{p}_L + M^2_{\tilde{m}_L} \tilde{m}_L^\dagger \tilde{m}_L + M'^2_{\tilde{\phi}} \tilde{\phi}^\dagger \tilde{\phi} + M'^2_{\tilde{\chi}} \tilde{\chi}^\dagger \tilde{\chi} \ , \tag{A.1}$$

where $M'^2_{\tilde{\phi}}$ and $M'^2_{\tilde{\chi}}$ are 2×2 hermitian matrices, $\tilde{\phi} \equiv (\tilde{\phi}_L, \tilde{\phi}_R)^\top$ and $\tilde{\chi} \equiv (\tilde{\chi}_L, \tilde{\chi}_R)^\top$. The mass terms in the superpotential also contribute to the $SU(2)_W$-singlet techniscalar mass matrices:

$$M^2_{\tilde{\phi}} = M'^2_{\tilde{\phi}} + m^2_\phi \check{I} \ , \tag{A.2}$$

where $\check{I}$ is the 2×2 unit matrix.

The squark and slepton mass terms are given by $(i, j = 1, 2, 3)$

$$\mathcal{L}'_{2s} = (M_L^{q\,2})_{ij} \tilde{q}_i^\dagger \tilde{q}_j + (M_R^{u\,2})_{ij} \tilde{u}_{i_R}^\dagger \tilde{u}_{jR} + (M_R^{d\,2})_{ij} \tilde{d}_{i_R}^\dagger \tilde{d}_{iR} + (M_L^{l\,2})_{ij} \tilde{l}_i^\dagger \tilde{l}_j + (M_R^{e\,2})_{ij} \tilde{e}_{i_R}^\dagger \tilde{e}_{jR} \ . \tag{A.3}$$

The 3×3 squark and slepton mass matrices are real, symmetric and have vanishing $(1, 2)$ and $(2, 1)$ elements (see Section 3).

The trilinear scalar terms in the super-weak basis, defined by the choice of the superpotential in Eq. (2.5) and of the squark and slepton mass terms in Eq. (A.3), are given by:

$$\begin{aligned} \mathcal{L}_{3s} \ = \ & \mu_{q_i} \epsilon^{\alpha\beta} \tilde{q}_{i_\alpha} \tilde{\Upsilon}_\beta \tilde{\phi}_L + \mu_{u_i} \tilde{u}_{i_R}^\dagger \tilde{m}_L \tilde{\phi}_R^\dagger + \mu_{d_i} \tilde{d}_{i_R}^\dagger \tilde{p}_L \tilde{\phi}_R^\dagger \\ & + \mu_{l_i} \epsilon^{\alpha\beta} \tilde{l}_{i_\alpha} \tilde{\Upsilon}_\beta \tilde{\chi}_L + \mu_{e_i} \tilde{e}_{i_R}^\dagger \tilde{p}_L \tilde{\chi}_R^\dagger + \text{h.c.} \ , \end{aligned} \tag{A.4}$$

where $\mu_{q_i}, \mu_{u_i}, \mu_{d_i}, \mu_{l_i}, \mu_{e_i}$, $i = 1, 2, 3$ are mass parameters. Although these terms are linear in squark and slepton fields, in the super-weak basis the flavors are uniquely defined and, in general, all the generations couple to the techniscalars.

Finally, the soft SUSY terms include majorana mass terms for all the gauginos: bino, winos, gluinos and techni-gluinos.

Appendix B

In this Appendix we derive the functions f^u_{ij} which enter in the elements of the up quark mass matrix (see Eq. (3.9)).

At energies higher than the technicolor scale $\Lambda \approx 4\pi v$ the technifermion condensate breaks and the chirality-flip mixing of the squarks vanishes. Hence, the integral corresponding to the one-loop graph in Fig. 2 should be cut off at Λ. The (i, j) $(i, j = 1, 2, 3)$ element of the up quark mass matrix is given by:

$$\begin{aligned} m^u_{ij} &= \frac{16}{3}\pi\alpha_s m_{\tilde{g}} m_\phi C_q C_t \langle \bar{m} m \rangle \\ &\times \int^\Lambda \frac{d^4p}{(2\pi)^4} \frac{-i}{(p^2 - m^2_{\tilde{g}})(p^2 - m^2_\phi)} \sum_{k,l=1}^{3} \frac{(U^q_L)^*_{3k}(U^q_L)_{ik}(U^u_R)^*_{jl}(U^u_R)_{3l}}{(p^2 - M^{q2}_{kL})(p^2 - M^{u2}_{lR})}\ , \end{aligned} \tag{B.1}$$

where M^{q2}_{iL} (M^{u2}_{iR}), are the eigenvalues of the left-handed (right-handed) up squark mass matrix, M^{q2}_L (M^{u2}_R), and U^q_L (U^u_R) is the unitary matrix which diagonalizes M^{q2}_L (M^{u2}_R, respectively). Integrating over the angles, using Eq. (3.4) and comparing Eqs. (3.9) and (B.1) we find

$$f^u_{ij} = \frac{4}{3}\int_0^{\Lambda^2/m^2_{\tilde{g}}} dy \frac{y}{(y+1)(y + m^2_\phi/m^2_{\tilde{g}})} \sum_{k,l=1}^{3} \frac{(U^q_L)^*_{3k}(U^q_L)_{ik}(U^u_R)^*_{jl}(U^u_R)_{3l}}{(y + M^{q2}_{kL}/m^2_{\tilde{g}})(y + M^{u2}_{lR}/m^2_{\tilde{g}})}\ . \tag{B.2}$$

The unitarity of U^q_L and U^u_R and the structure of the integrand in Eq. (B.2) indicates an upper bound $f^u_{ij} \lesssim \mathcal{O}(10^{-1})$.

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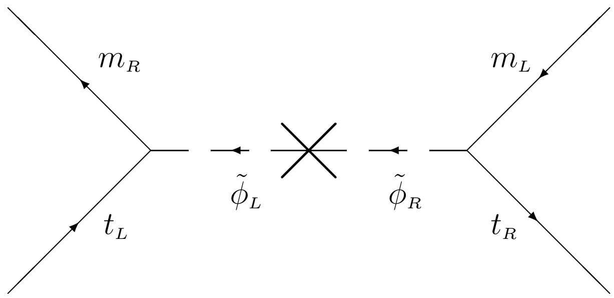
Figure captions

Fig. 1. Four-fermion interaction due to the exchange of technicolored scalars. The cross on the scalar line denotes the chirality-flip mixing of the $\tilde{\phi}$ scalars.

Fig. 2. The radiative correction involving a gluino, $\tilde{g}$, gives the leading contribution to the masses of the up, u_1, and charm, u_2, quarks.

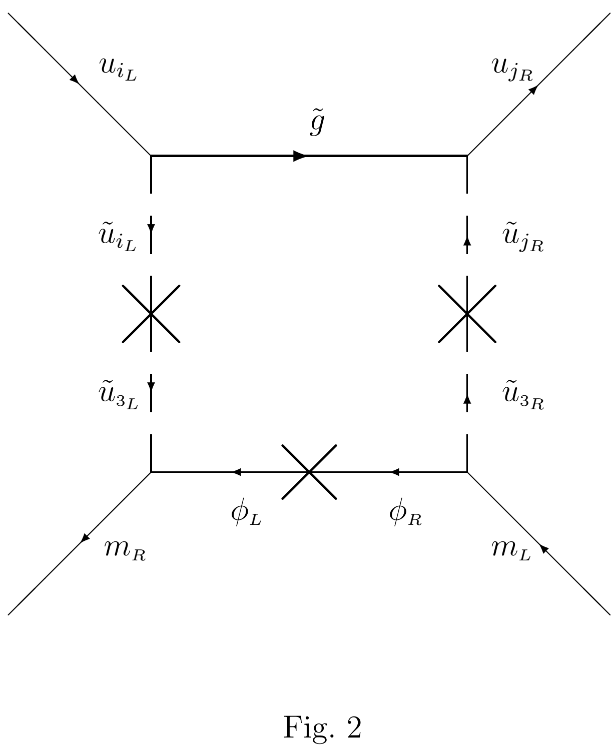
Fig. 3. The leading contribution to m_d has a techni-gluino, $\tilde{\mathcal{G}}_{TC}$, and techniscalars in the loops.

Fig. 4. Technifermion exchange leading to a dimension-five operator responsible for $\tilde{d}_{3_L}$-$\tilde{d}_{3_R}$ mixing. The cross on the fermion line represents the $SU(2)_W \times U(1)_Y$ singlet mass term for the ϕ technifermion.



m_R
m_L
$\tilde{\phi}_L$
$\tilde{\phi}_R$
t_L
t_R

Fig. 1



u_{i_L}
u_{j_R}
$\tilde{g}$
$\tilde{u}_{i_L}$
$\tilde{u}_{j_R}$
$\tilde{u}_{3_L}$
$\tilde{u}_{3_R}$
ϕ_L
ϕ_R
m_R
m_L

Fig. 2

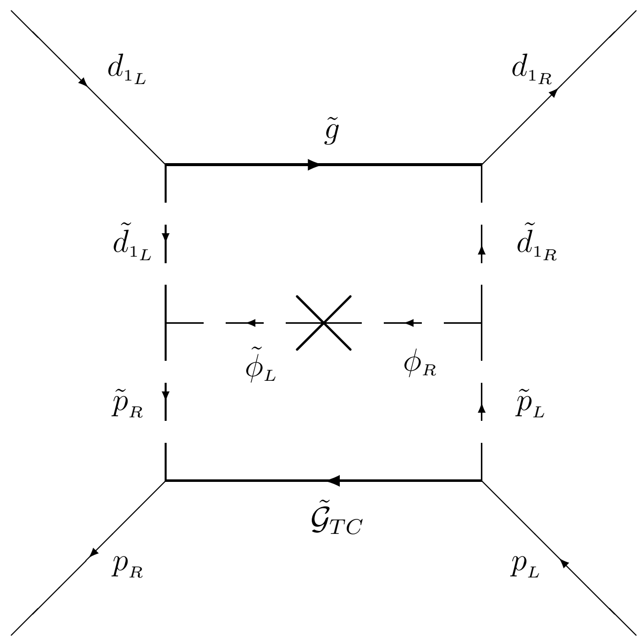

d_{1_L}
d_{1_R}
$\tilde{g}$
$\tilde{d}_{1_L}$
$\tilde{d}_{1_R}$
$\tilde{\phi}_L$
ϕ_R
$\tilde{p}_R$
$\tilde{p}_L$
$\tilde{\mathcal{G}}_{TC}$
p_R
p_L

Fig. 3

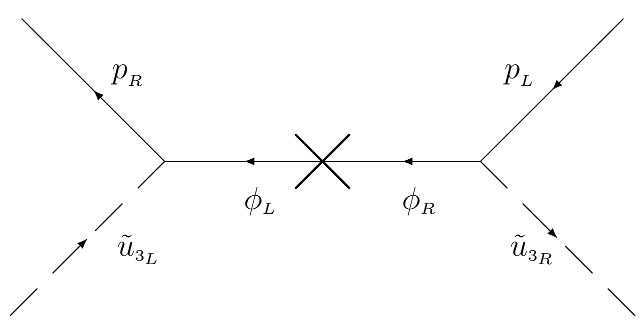

p_R
p_L
ϕ_L
ϕ_R
$\tilde{u}_{3_L}$
$\tilde{u}_{3_R}$

Fig. 4